



On Quarks and Flavour Symmetry

H. Hogaasen, Paul Sorba

► To cite this version:

| H. Hogaasen, Paul Sorba. On Quarks and Flavour Symmetry. 2007. hal-00189240

HAL Id: hal-00189240

<https://hal.science/hal-00189240>

Preprint submitted on 20 Nov 2007

HAL is a multi-disciplinary open access archive for the deposit and dissemination of scientific research documents, whether they are published or not. The documents may come from teaching and research institutions in France or abroad, or from public or private research centers.

L'archive ouverte pluridisciplinaire **HAL**, est destinée au dépôt et à la diffusion de documents scientifiques de niveau recherche, publiés ou non, émanant des établissements d'enseignement et de recherche français ou étrangers, des laboratoires publics ou privés.

On Quarks and Flavour Symmetry

H. Høgaasen

*Department of Physics University of Oslo
Box 1048 NO-0316 Oslo Norway
hallstein.hogasen@fys.uio.no*

and

P. Sorba

*LAPTH¹, Université de Savoie, CNRS
9 chemin de Bellevue, B.P. 110, F-74941 Annecy-le-Vieux Cedex,
FRANCE
sorba@lapp.in2p3.fr*

Abstract

Hadronic spectroscopy can be introduced to students and developed rather far without requiring $SU(N)$ flavour symmetry. In such a "minimalist" presentation, we are naturally led to comment and clarify the concept of the "generalized" Pauli principle.

¹Laboratoire d'Annecy-le-Vieux de Physique Théorique, UMR 5108

1 Introduction

Internal symmetry groups have a glorious place in the history of physics. One of the highlights was the discovery by Fermi that low energy pion-nucleon scattering is dominated by a single resonance of spin $3/2$. By treating the nucleons as an isospin doublet, the pions as an isospin triplet and the resonance Δ as a state with isospin $3/2$, Fermi and his collaborators proved the dynamical fact that, to a good precision, isospin is conserved in strong interactions.

Another highlight was Gell-Mann's prediction of the spin $3/2$ Ω^- at the Geneva conference in 1962. The subsequent discovery of this state at Brookhaven in 1964 made everyone confident that the *flavour* – group $SU(3)$ was as relevant for strong interactions as the isospin group.

There is however a difference between the two cases we have mentioned: Fermi's discovery was a discovery in the dynamics of particles while Gell-Mann's prediction is usually presented as coming from the assumption that baryons should fall into specific representations of $SU(3)$ -flavour, namely the 8 dimensional one for the lightest spin $1/2$ baryons, the 10 dimensional one for the lightest spin $3/2$ baryons.

Experimental results from reactions involving the assignment of mesons and baryons to irreducible representations of $SU(3)$ came later, and showed that the concept of "broken flavour symmetry" was useful.

Subsequently the flavour symmetry groups $SU(4)$ and $SU(5)$ have been invoked to classify the multiplicities of states when charm and bottom quantum numbers are added.

But let us stress that the assignment of a multiplet of particles to a group representation is in itself an empty statement. Its usefulness depends on the group to be an (almost) symmetry group. When Fermi's team discovered that the isospin group - historically, this group was first supposed to be $R(3)$ before becoming $SU(2)$ - it was by studying reactions. If a pion-nucleon state is a linear combination of isospin $I = 1/2$ and $I = 3/2$, only two independent amplitudes $A_{1/2}$ and $A_{3/2}$ will describe all $\pi N \rightarrow \pi N$ reactions ²: this

²As $A_{3/2}$ dominates at low energy, the analysis of data were quite simple.

turned out to be correct.

Experiments must decide to which extent flavour symmetry is a useful concept. From the multiplication table of SU(3):

$$8 \otimes 8 = 1 \oplus 8 \oplus 8 \oplus 10 \oplus \overline{10} \oplus 27$$

it follows that the 64 reactions we have in "quasi elastic" $2 \leftrightarrow 2$ reactions between octet states are given by only six independent amplitudes. From the same table we see, as the representation 10 occurs only once, that in the coupling of a decuplet to two octets there is only one amplitude. This is similar to the decay of $\Delta \rightarrow N\pi$ in the isospin symmetric - SU(2) - case. So, if one has baryons made from u, d and s quarks, the assumption of flavour symmetry relates the coupling constant in $\Delta^{++} \rightarrow P\pi^+$ to the coupling constants in $\Sigma^{*0} \rightarrow \Lambda\pi^0$, $\Sigma^{*0} \rightarrow \Sigma^+\pi^-$ and $\Xi^{*0} \rightarrow \Xi^0\pi^0$. Experimentally they come out correctly within 10% .

1964 was the year Gell-Mann and Zweig invented quarks.

Right from the start Zweig was a fervent believer in the objective existence of quarks (that he called aces) as the fundamental constituents of hadrons. With three flavours of aces (quarks) mesons came in $q\bar{q}$ nonets. There was no need to invoke a "magic mixing" between a flavour octet and a singlet; to explain that one of the vector mesons decayed into $K\bar{K}$ pairs was simply the reflection of the fact that it was composed of $s\bar{s}$. If some "magic" is there , it is rather in the structure of the spin zero mesons.

In the years before quarks were generally accepted as physical beings in the physics community, flavour symmetry arguments were (and often still are!) regarded as more high-brow than arguments simply based on quarks. But today hardly anybody denies that quarks exist. It is easy to criticize the (nonrelativistic) quark model where baryons are made from three valence quarks and mesons from a quark-antiquark pair, but it cannot be denied that it gives a set of rules to classify hadron states, and also to calculate many of their static properties which are quite illuminating.

However, as we will see in the following, we do not need any flavour $SU(N)$ group either to classify hadrons or to estimate many of their physical properties. Hereafter, the fundamental symmetry groups we will use are the group of rotations in the three dimensional space represented by the $SU(2)$ spin group (together with the $O(3)$ orbital angular momentum one in the case of excited states), and the colour $SU(3)$ group.

Indeed, the $SU(3)$ symmetry group of fundamental importance is not the group acting over flavours, but the group acting over colour space. It is a group that is gauged and thereby introduces gluonfields while defining QCD. The (yet unproven) dogma that free particles are colourless, i.e. transform as a singlet under the colour $SU(3)$ group, implies that the most economical configurations are those involving three quarks (baryon) or a quark-antiquark pair (meson). This comes from the decomposition of (colour) $SU(3)$ representations:

$$3 \otimes 3 \otimes 3 = 1 \oplus 8 \oplus 8 \oplus 10 \quad \text{and} \quad 3 \otimes 3 = 6 \oplus \bar{3} \quad 3 \otimes \bar{3} = 1 \oplus 8$$

where the representations 6 and $\bar{3}$ respectively involve symmetric and anti-symmetric combinations of the two colour indices.

The second essential ingredient we need is the Pauli principle. Of course it is of importance only for baryons³, and it is actually this class of hadrons which deserves to be considered for our purpose. Indeed, at least with regard to their classification, $q\bar{q}$ mesons with a given spin and parity are simply gathered in multiplets of dimension N^2 , N being the number of different quark flavours.

Note that such three quark states are colour-antisymmetric under the exchange of two quarks. Then, the Pauli principle tells us that a state involving two identical quarks, i.e. quarks with identical flavours, must be symmetric in the other quantum numbers.

We will start by showing in Sections 2 and 3 that the multiplets of low-lying as well as excited baryonic resonances can be directly constructed

³We do not consider mixing with multi-quark states. For mesons we need a generalization of the Pauli principle which follows from local field theory through the CPT operation; but this point will not be discussed hereafter where we limit to (usual) baryons.

with the above hypotheses and tools at hand. In Section 4, expressions for S-wave three quark states only satisfying permutational symmetry imposed by the Pauli principle are given and used to compute baryon mass splittings through (colour) magnetic interactions between the constituent quarks; their magnetic moments are also obtained.

However the basis constructed in Section 2 does not appear suitable for computations of physical quantities involving flavour changing forces, such as, for instance, semi-leptonic decays. A convenient way to circumvent this difficulty is to impose complete symmetry under permutations of the three constituents within the baryon wave function rather than the - partial-symmetry dictated by the Pauli principle: this prescription appeals to the so-called "generalized" Pauli principle, which, as is well known, does not introduce any extra physical assumption. The corresponding developments are presented in Section 5.

We conclude by commenting on the direct connection between irreducible representations of the unitary groups and their properties under the permutation group, which explains, in our opinion, their perfect adaptability to the classification of baryons.

In order not to overload the text and keep clear the main idea of this note, we have decided to treat in appendices the following points. Appendix A is devoted to a general and explicit construction of the (three state particle) Jacobi coordinates. In our knowledge, such coordinates have already been explicated in the case of the harmonic oscillator with only one coupling constant, but not with a different coupling associated to each couple of constituents.

Then we give a proof, in Appendix B, of the equivalence between the two Hilbert space descriptions introduced in Section 4 and 5 respectively, which involve states satisfying partial permutation symmetry (simple Pauli principle) for the former, and total permutation symmetry (generalized Pauli principle) for the later. The Fock space emerges rather naturally from this construction, and we show in Appendix C how it can be used to represent baryon states. Finally, we give in Appendix D the explicit wave-functions of the eight $S = 1/2$ low-lying baryons in the three bases successively in-

troduced: the permutational "partial" symmetric basis (cf. Section 4), the totally symmetric basis (cf. Section 5) and the Fock space basis, this last one combining the advantages of the other two representations: simplicity of the expressions and complete symmetry.

As this paper can be partly regarded as a set of lecture notes, there is an almost empty reference list. We apologize for this but we realize that to do justice to all the people who developed the subject, we would need a list longer than the article.

2 The ground state of three quarks

This is the subject of most introductory courses on elementary particles. So, we will treat this first. Although the following is known, it is apparently not well known, as we realized by looking up at lectures notes on the web. Among many popular monographs, we found only one [1] which shares the approach presented below.

We now forget about colour, remembering that each pair of quarks are antisymmetric in colour so that the baryon states must be symmetric in the other degrees of freedom. This was for a long time called the Dalitz symmetric quark model. Note at this point that it was realized [2] that if quarks were fermions such that there existed an hidden quantum number - which later turned out to be the colour - justifying the symmetric quark model, then the multiplicity of the ground states would follow independently of any flavour symmetry group.

In the ground state where no angular momentum is involved, the only degree freedom of each quark is the spin. As a pair of identical quarks must be symmetric under interchange it must be symmetric in spin: the spin of the pair is one.

Integrating out the spatial degrees of freedom, we are left with a Hamiltonian over the flavour-spin space of the quarks.⁴ Now let us count the ground states:

Suppose that we have N flavours (u, d, s, c,...) and we choose three of these to make baryons.

⁴The following arguments are not restricted to nonrelativistic quantum mechanics.

First for three identical flavours:

There are N ways of choosing these baryons made up of three identical quarks. Each pair has $S=1$ so all have total spin $3/2$.

Two identical flavours:

There are $N(N-1)$ ways of choosing two identical and one different from the two first chosen.

The identical quarks are coupled to spin one, and the total spin is $S=3/2$ or $S=1/2$.

All three quarks different:

There are $N(N-1)(N-2)/6$ ways of choosing three different flavours. The Pauli principle imposes no restriction here, so the total spin can be $3/2$ or $1/2$, the multiplicity of $S=1/2$ is two.

So the number $N(3/2)$ of flavour states with total spin $3/2$ is:

$$N(3/2) = N + N(N-1) + N(N-1)(N-2)/6 = N(N+1)(N+2)/6$$

and the number of flavour states with spin $1/2$ is:

$$N(1/2) = N(N-1) + 2N(N-1)(N-2)/6 = N(N+1)(N-1)/3$$

We immediately see that we have found the same multiplicities as is commonly inferred from the dimensions of the representations of the flavour symmetry group $SU(N)$.

For $N=3$ we have an "eightfold way", more precisely an octet of $S=1/2$ states and a decuplet of $S=3/2$ states.

But we also realize that these multiplicities have nothing to do with the existence of an internal symmetry group, they have their sole origin in the Pauli principle. Not only u, d, s , quarks, but also u, c, b , or any triplet of flavour quarks will provide with an octet and a decuplet.

Now, let us count the total number of such quantum mechanical states. Since there are $2S+1$ states in a spin S representation, one gets:

$$N(\text{total}) = 4 N(N+1)(N+2)/6 + 2 N(N+1)(N-1)/3 = 2 N(N+1)(2N+1)/3$$

and that is exactly the dimension of the $SU(2N)$ completely symmetric representation arising from the tensorial product of three $2N$ dimensional fundamental $SU(2N)$ representations. Taking as an example $N=3$, we indeed get

56 states, that is the dimension of the corresponding symmetric representation of the usually called "flavour-spin" SU(6) group: we will comment later on this point, in direct connection with the "generalized" Pauli principle.

3 Excited states

Let us now turn our attention to the construction of P-wave baryon states, to which we will restrict our presentation in order not to overload this note. These negative parity baryons are associated to an L=1 orbital momentum. This unit of O(3) orbital angular momentum stands usually in one of the two quark relative coordinates, and the most common model which is well adapted to represent such an effect is that of the harmonic oscillator. The reason for treating confining forces between quarks using harmonic oscillator potentials is simply that the center of mass motion can be separated out using Jacobi coordinates. Note that this choice has no importance for our purpose, namely the *number* of excited states.

We will use as a Hamiltonian

$$H_{\text{h.o.}} = \sum_i (m_i + p_i^2/2m_i) + \sum_{i<j} (k/2) \cdot (\vec{r}_i - \vec{r}_j)^2 \quad (1)$$

where \vec{r}_i , $i=1,2,3$, denote the respective positions of the three quarks. Although we do not start with a more general Hamiltonian H_{gen} , we can obtain the general solution for the spectrum by first using $H_{\text{h.o.}}$ to get states of a harmonic oscillator basis and then perturb them with the perturbation $H_{\text{gen}} - H_{\text{h.o.}}$. The energy of the levels will change with respect to the ones of the harmonic oscillator, but their number will be the same.

In terms of the Jacobi coordinates $\vec{\rho}$, $\vec{\lambda}$ and the center of mass coordinates \vec{R}_{cm} , the Hamiltonian separates into $H_{\text{h.o.}} = \sum_i m_i + H_{\text{cm}} + H_{\rho} + H_{\lambda}$ with

$$\begin{aligned} H_{\text{cm}} &= P^2/2M, \quad M = \sum_i m_i, \quad \vec{P} = \sum_i \vec{p}_i, \\ H_{\rho} &= p_{\rho}^2/2m_{+} + 3k/2 \cdot (\vec{\rho})^2 \\ H_{\lambda} &= p_{\lambda}^2/2m_{\lambda} + 3k/2 \cdot (\vec{\lambda})^2. \end{aligned} \quad (2)$$

The explicit expressions of m_{+} and m_{λ} are given in Appendix A, where a detailed construction of Jacobi coordinates is developed, and also generalized

to the case where we have a different coupling k_{ij} for each quark pair $q_i q_j$ in the potential term.

For notational simplicity, we consider here the equal mass case, i.e. $m_i = m$ to which correspond: $m_+ = m_- = m$ and the relative coordinates $\vec{\rho}$ and $\vec{\lambda}$ reduce to :

$$\begin{aligned}\vec{\rho} &= (\vec{r}_1 - \vec{r}_2)/\sqrt{2} \\ \vec{\lambda} &= (\vec{r}_1 + \vec{r}_2 - 2\vec{r}_3)/\sqrt{6}\end{aligned}$$

So, let us separately study the different configurations:

i) **Baryons made of three flavour identical quarks qqq :**

Consider the oscillator relative to $\vec{\rho}$: it is antisymmetric in the exchange (1) -(2). The doublet made with the first two quarks qq , that is of quarks in position (1) and (2), has automatically spin $S = 0$, in order for the total spin and orbital momentum part to be symmetric (Pauli principle). It follows that in this $\vec{\rho}$ configuration, the total spin of the baryon is $S = 1/2$, and the spin/orbit part of the wave function reads, for $S_z = +1/2$ and up to a normalization factor:

$$(\vec{r}_1 - \vec{r}_2) (\uparrow\downarrow\uparrow - \downarrow\uparrow\uparrow) + \text{sym.}$$

Replacing now $\vec{\rho}$ by $\vec{\lambda}$, which is symmetric in the first two quarks, one easily deduces that the corresponding qq doublet has a (symmetric) spin $S = 1$. However in the product $S = 1$ by $S = 1/2$, the totally spin symmetric $S = 3/2$ part cannot provide, when combined with the (not completely symmetric) $\vec{\lambda}$, a completely symmetric wave function. It follows that the only possibility for the resulting baryon is to have spin $1/2$. Moreover, one notes that the spin/orbit part of the wave function reads:

$$(\vec{r}_1 + \vec{r}_2 - 2\vec{r}_3) (2\uparrow\uparrow\downarrow - \downarrow\uparrow\uparrow - \uparrow\downarrow\uparrow) + \text{sym.}$$

which is exactly the same, up to a scale factor, as the one obtained just above for the $\vec{\rho}$ case.

Therefore, for qqq configurations, the only solution is given by the $\vec{\rho}$ oscillator with two quark spin one *or* by the $\vec{\lambda}$ oscillator with two quark spin zero, and total spin $S = 1/2$.

ii) **Baryons made of two (and only two) identical flavour quarks qqq' :**

Now the Pauli principle imposes permutation symmetry only in the first two quarks qq . Note that, as explicated in Appendix A, the antisymmetry (resp. symmetry) of \vec{p} (resp. $\vec{\lambda}$) persists when $m_1 = m_2 = m$ and $m_3 = m'$ with m and m' different. Then one gets:

- with the \vec{p} oscillator: the qq spin part must be zero, leading for the (qqq') baryon to $S = 1/2$, with in the wave function a qq spin/orbit part proportional to:

$$(\vec{r}_1 - \vec{r}_2) (\uparrow\downarrow - \downarrow\uparrow) + \text{sym.}$$

- with the $\vec{\lambda}$ oscillator: the qq spin part is now one, allowing the total spin to be $S = 3/2$ and $S = 1/2$.

iii) **Baryons made of three different flavour quarks $qq'q''$:**

Now there is no restriction imposed by the Pauli principle. It follows that with the \vec{p} as well as with the $\vec{\lambda}$ oscillator, one can get one baryon with total spin $S = 3/2$ and two baryons of total spin $S = 1/2$ for each triplet of flavours.

As a conclusion, let us count the number of states of three different flavours:

- with $S = 3/2$: there are 6 states of qqq' -type with the $\vec{\lambda}$ oscillator, one state $qq'q''$ with each of the \vec{p} and $\vec{\lambda}$ oscillator, that is a total of 8 states.

- with $S = 1/2$: there are 3 states of qqq -type made from the \vec{p} osc. (and partial qq spin $S = 0$) or from the $\vec{\lambda}$ osc. (and partial qq spin $S = 1$). There are also 6 states of qqq' -type made from the \vec{p} osc. (and partial qq spin $S = 0$), 6 states of qqq' -type made from the $\vec{\lambda}$ osc. (and partial qq spin $S = 1$). Finally, there are 4 states of $qq'q''$ -type respectively made with the \vec{p} and $\vec{\lambda}$ osc. and of qq partial spin $S = 0$ and 1.

Then, with three flavours, one obtains a total of 70 states, with 8 states of $S = 3/2$ and 19 states of $S = 1/2$ in perfect accordance with the usual SU(6) approach where the P-wave baryons ($L = 1$) are classified in the irreducible 70 dimensional SU(6) representation, itself decomposing with respect to SU(3) flavour and SU(2) spin as:

$$70 = (8, 3/2) + (8 + 10 + 1, 1/2)$$

In a similar manner, one can show that, for higher excitations, the number of excited states and their configurations are always what one should infer from the SU(6) approach.

4 Flavour nonchanging forces

Let us now focus on forces acting on the quarks (gluonic and electromagnetic in particular) that do not change the quark flavour.

We can represent any baryon made of the three quarks q_1, q_2, q_3 as follows:

$$B(q_1, q_2, q_3) = (q_1 q_2)_s \otimes (q_3)_{1/2} \quad (3)$$

where s is the spin of the doublet of the quark pair $q_1 q_2$. One obviously has $S = 1$ or $S = 0$.

If we denote by q_i^\uparrow the i quark with spin up, by q_i^\downarrow the corresponding spin down state, then with the help of Clebsch-Gordan coefficients, any $S = 3/2$ state with $S_z = 3/2$ writes :

$$q_1^\uparrow q_2^\uparrow q_3^\uparrow$$

while it becomes for $S_z = 1/2$:

$$[q_1^\uparrow q_2^\uparrow q_3^\downarrow + q_1^\uparrow q_2^\downarrow q_3^\uparrow + q_1^\downarrow q_2^\uparrow q_3^\uparrow]/\sqrt{3}$$

and so on. As we have seen there are 10 such flavour states, if we limit ourselves to three flavours.

For the 8 flavour states with spin 1/2 and $S_z = +1/2$, one must distinguish the case when the first two quarks are identical in flavour from the case when all flavours are different. In the first situation, the spin of q_1 and q_2 must couple to one (we write below: $q_1 = q_2$) and Clebsch-Gordan coefficients give:

$$\frac{1}{\sqrt{6}} [2 \cdot q_1^\uparrow q_1^\uparrow q_3^\downarrow - (q_1^\uparrow q_1^\downarrow + q_1^\downarrow q_1^\uparrow) q_3^\uparrow]$$

. In the second configuration, the Pauli principle gives no restriction, and we have two states, the spin of the two first quarks being coupled to produce either a spin one (ψ_1) or zero (ψ_2):

$$\psi_1 = \frac{1}{\sqrt{6}} [2 \cdot q_1^\uparrow q_2^\uparrow q_3^\downarrow - (q_1^\uparrow q_2^\downarrow + q_1^\downarrow q_2^\uparrow) q_3^\uparrow]$$

$$\psi_2 = [(q_1^\uparrow q_2^\downarrow - q_1^\downarrow q_2^\uparrow)q_3^\uparrow]/\sqrt{2}$$

It clearly does not matter which ordering we have for the different flavours, but it will be shown in a moment that it is often convenient to order them by placing the lightest quark(s) in front of the heaviest.

We stress again that the *classification* of states is into a spin 3/2 decuplet and a spin 1/2 octet, *no matter* which three flavours we choose. With quarks d, c and b we have the same structure as with the lightest states u, d and s. From this remark it should be clear for students that the classification of states is one thing; the question whether we have a flavour symmetry group SU(3) is completely different.

Now, from the 8 states with total spin 1/2 that we have constructed, we can span a Hilbert space and consider how interactions act on states therein.

If there were no spin dependent interquark forces, it is evident that the states made of the same quarks would have the same mass. The interquark Hamiltonian would be diagonal in our Hilbert space.

This is definitely not so. In the beginning it seemed strange that the Λ and Σ_0 had different masses. An early explanation, based on spin-spin interactions, was given by Sakharov and Zeldowich in 1966 with the Hamiltonian :

$$H_S = \sum_{i,j} c_{ij} \vec{\sigma}_i \cdot \vec{\sigma}_j \quad (4)$$

all coefficients c_{ij} being equal. In 1975 this spin-spin interaction was shown to be a consequence of QCD by De Rujula, Georgi and Glashow:

$$H_{CM} = - \sum_{i,j} C_{ij} \vec{\lambda}_i \cdot \vec{\lambda}_j \vec{\sigma}_i \cdot \vec{\sigma}_j \quad (5)$$

As all quark pairs are in a $\bar{3}$ state of colour, the colour part $\vec{\lambda}_i \cdot \vec{\lambda}_j$ factorizes, giving a common factor - 8/3, so that:

$$H_{CM} = (8/3) \sum_{i,j} C_{ij} \vec{\sigma}_i \cdot \vec{\sigma}_j \quad (6)$$

Here the coefficients C_{ij} are, among other things, dependent on the quark masses and properties of the spatial wave functions of the quarks in the

system.

A natural scaling assumption for the coefficients $C_{ij} \propto \frac{1}{m_i m_j}$ comes from the analogy with the hyperfine splitting in atoms that originates in the interaction between (electro) magnetic moments. In the physics of quarks the analog interaction is between the quarks *colourmagnetic* moments.

Let us now imagine that we have integrated out all the spatial variables for the three quarks. We can then write an effective Hamiltonian over the spin-space of the quark:

$$H = \sum_i m_i + (8/3) \sum_{i,j} C_{ij} \vec{\sigma}_i \cdot \vec{\sigma}_j \quad (7)$$

Here effective masses m_i incorporate the masses of the quarks as well as their kinetic energy⁵. If we *assume* that each effective mass m_i is (almost) the same in different baryons, then we get mass formulae. For this, we have only to determine the eigenvalues of H_{CM} .

The solution of the eigenvalue problem comes easily when one uses the following identity for the Pauli matrices - as can be directly tested out by applying $\vec{\sigma}_i \cdot \vec{\sigma}_j$ on the (symmetric) spin one state and on the (antisymmetric) spin zero state:

$$\vec{\sigma}_i \cdot \vec{\sigma}_j = -1 + 2P_{i \leftrightarrow j} \quad (8)$$

where $P_{i \leftrightarrow j}$ is the operator that permutes the spin states of the two particles i and j. One sees at once that the Hamiltonian is almost diagonal in our Hilbert space.

For states with total spin $S = 3/2$ the eigenvalue of H_{CM} is then:

$$\frac{8}{3}(C_{12} + C_{13} + C_{23}) \quad (9)$$

while for the spin $S = 1/2$ baryons, it reduces to :

$$\frac{8}{3}(C_{12} - 4C_{13}) \quad (10)$$

⁵If we added on $H_S = -\sum_{i,j} c_{ij} \vec{\sigma}_i \cdot \vec{\sigma}_j$ we would get the most general Hamiltonian we can have for the system of three quarks when the spatial variable are integrated out. So H is more general than *one* gluon exchange only.

for all states with two identical quarks (they are chosen above as being q_1 and q_2). The mass of corresponding $S=1/2$ states reads then:

$$2m_1 + m_3 + \frac{8}{3}(C_{12} - 4C_{13}) \quad (11)$$

The only two states that are mixed are the spin $1/2$ states where all three flavours are different. Using the same spin coupling scheme as before, where the spin of the first two quarks is coupled to one (ψ_1) or zero (ψ_2), and where the total spin is $1/2$:

$$\begin{aligned} \psi_1 &= |(q_1 q_2)_1\rangle \otimes |(q_3)_{1/2}\rangle \\ \psi_2 &= |(q_1 q_2)_0\rangle \otimes |(q_3)_{1/2}\rangle \end{aligned} \quad (12)$$

one easily finds the colour-spin Hamiltonian over these two states as:

$$H_{\text{CM}} = \frac{8}{3} \begin{bmatrix} C_{12} - 2(C_{13} + C_{23}) & -\sqrt{3}(C_{13} - C_{23}) \\ -\sqrt{3}(C_{13} - C_{23}) & -3C_{12} \end{bmatrix} \quad (13)$$

Now we see that if $C_{13} = C_{23}$ this matrix is also diagonal. If q_1 is the u -quark and q_2 is the d -quark, we would expect this to be approximately true. The same remark holds for the effective masses $m_1 \simeq m_2$. Then we note that in this approximation the largest eigenvalue leads to a mass:

$$2m_1 + m_3 + \frac{8}{3}(C_{12} - 4C_{13}) \quad (14)$$

and that is the mass of both the uuq_3 and ddq_3 systems given by expression (11). Therefore we have *three* states in the octet made from quarks u , d and q_3 (q_3 being s , c or b), with the same mass (as they would have if we had an isospin 1 state and isospin invariance in nature) .

The approximate isospin independence that sits in the QCD Lagrangian - due to the smallness of the (current) u and d mass compared to Λ_{QCD} , - reappears in the masses of the three-quark bound states.

If we make the baryons out of the lightest quarks u , d and s , we see that the states we called ψ_1 and ψ_2 are those commonly denoted Σ_0 and Λ .

The mixing between ψ_1 and ψ_2 - in this case called Λ - Σ_0 mixing - is induced by an (isospin breaking) inequality $C_{13} \neq C_{23}$.

So this type of mixing is the general one for all ψ_1 and ψ_2 states, it is quite

small for uds , udc , udb , much bigger for states like ucb .

In the general case the eigenvectors for H_{cm} are :

$$V_+ = \cos\theta \cdot \psi_1 + \sin\theta \cdot \psi_2$$

$$V_- = -\sin\theta \cdot \psi_1 + \cos\theta \cdot \psi_2$$

with corresponding eigenvalues for H_{cm} :

$$\lambda_{\pm} = \frac{8}{3} [-(C_{12} + C_{13} + C_{23}) \pm 2 \sqrt{C_{12}^2 - C_{12} C_{13} - C_{12} C_{23} + C_{13}^2 - C_{13} C_{23} + C_{23}^2}]$$

and mixing angle θ given by

$$\tan(2 \cdot \theta) = \frac{-\sqrt{3} \cdot (C_{13} - C_{23})}{2 \cdot C_{12} - (C_{13} + C_{23})} \quad (15)$$

For light ($q = u, d$ and s) quarks, baryon masses are well reproduced with masses

$$m_q = 360, \quad m_s = 535 \text{ MeV}, \quad (16)$$

and strength factors

$$C_{qq} = 18.5, \quad C_{qs} = 12.5, \quad C_{ss} = 9.5 \text{ MeV}. \quad (17)$$

If one accepts that effective masses could differ when heavy quarks are present, a possible choice for S-wave charmed baryon states could be

$$m_c = 1550, \quad m_q = 450, \quad m_s = 590 \text{ MeV}, \quad (18)$$

with associated strength factors

$$\begin{aligned} C_{qq} &= 20, & C_{qc} &= 5, & C_{qs} &= 15, \\ C_{ss} &= 10, & C_{cs} &= 4, & C_{cc} &= 4 \text{ MeV}. \end{aligned} \quad (19)$$

The states we have constructed are considerably simpler than those one often ⁶ finds in textbooks and lecture notes. No mention is here of "mixed symmetry octet states" or "generalized Pauli principle". The states are symmetrized only where the Pauli principle demands it, and can be used for all calculations where flavour is not changed.

⁶But not always !

If one wants to complicate calculations by symmetrizing in all three flavours, one is of course free to do so ⁷.

Mathematically there is *indeed* a one to one correspondence between the states that we have used and the states that are symmetrized in all three particles. But to ask a student to compute the magnetic moment of the nucleons by using nine terms in the state, instead of three, is not very kind. That calculation is very simple indeed, using the magnetic moment operator for a baryon:

$$\vec{\mu} = \sum_i Q_i \cdot \mu_i \cdot \vec{\sigma}_i$$

where in the definition of μ_i we have taken out the electric charge factor Q_i of the quark i , but not the expected $\frac{1}{m_i \cdot c}$ dependence. All the μ_i 's are therefore positive.

The expectation value of $\vec{\mu}$ for a $|B, S = 1/2, S_z = 1/2\rangle$ state, composed of two identical quarks α and a different one β gives the value of the magnetic moment of the baryon B:

$$\mu_B = \langle B, 1/2, 1/2 | \mu_z | B, 1/2, 1/2 \rangle = \frac{4}{3} \cdot Q_\alpha \cdot \mu_\alpha - \frac{1}{3} \cdot Q_\beta \cdot \mu_\beta$$

With these conventions one gets for the proton and neutron: $\mu_P = \frac{8}{9}\mu_u + \frac{1}{9}\mu_d$, $\mu_N = -\frac{2}{9}\mu_u - \frac{4}{9}\mu_d$ and so on.

In the case where all three flavours are different, $q_1 = \alpha$, $q_2 = \beta$ and $q_3 = \gamma$, one finds for the two states

$$\begin{aligned} \mu_{\psi_1} &= \langle \psi_1 | \mu_z | \psi_1 \rangle = (2 \cdot Q_\alpha \cdot \mu_\alpha + 2 \cdot Q_\beta \cdot \mu_\beta - Q_\gamma \cdot \mu_\gamma) / 3 \\ \mu_{\psi_2} &= \langle \psi_2 | \mu_z | \psi_2 \rangle = Q_\gamma \cdot \mu_\gamma \end{aligned}$$

We include the off-diagonal matrix element for those that want to compute the effect of ψ_1 ψ_2 mixing:

$$\langle \psi_2 | \mu_z | \psi_1 \rangle = \langle \psi_1 | \mu_z | \psi_2 \rangle = (-Q_\alpha \cdot \mu_\alpha + Q_\beta \cdot \mu_\beta) / \sqrt{3}$$

Neglecting small mixing, the magnetic moment of Λ is given as $\mu_\Lambda = -\frac{1}{3}\mu_s$. A tolerable estimation of all the magnetic moments of the lightest baryons

⁷In this case our notation is not optimal. The coefficients C_{ij} should then be labeled by the quark flavours and not by numbers.

can be obtained by using as input the observed nucleon and Λ values, i.e.

$$\mu_u = \mu_d = \mu_P, \mu_s = -3 \cdot \mu_\Lambda.$$

The student will remark that $\mu_s < \mu_u$ as expected, but also that not all is well with this kind of calculation. An obvious problem is that experimentally $\mu_\Lambda - \mu_{\Xi^-}$ is positive, whereas this conventional type of calculation gives $\mu_\Lambda - \mu_{\Xi^-} = (\mu_s - \mu_d)/9$ which should be negative!

We now turn to interactions that change flavour, and we shall see that permutation symmetry reestablishes something equivalent to the "generalized Pauli principle" without invoking any "principle", just the consistency of the Hilbert space.

5 Flavour changing forces and the emergence of the generalized Pauli principle.

Suppose now that we want to teach flavour changing weak decays, as it is the case in semi-leptonic hyperon decays. Let us take the example of an s-quark turning into a u-quark as in the $\Lambda \rightarrow P e^- \bar{\nu}$ decay.

The state Λ is constituted by three quarks of different flavours: u, d, and s, so the Pauli principle does not bring any constraint in this case. This resonance has a total spin $S = 1/2$, and its (ud) part is of spin $S = 0$ when we ignore $\Lambda\Sigma_0$ mixing. We are therefore inclined to write it as:

$$\Lambda = (u^\uparrow d^\downarrow - u^\downarrow d^\uparrow) s^\uparrow / \sqrt{2}$$

In the transformation $s \rightarrow u$, the above expression changes into:

$$(u^\uparrow d^\downarrow - u^\downarrow d^\uparrow) u^\uparrow / \sqrt{2}$$

and this state is not in the Hilbert space defined in the previous section! It is not a state allowed by the Pauli principle.

The result is a state vector where the first term is symmetric under the interchange of the two u-quarks -and that is fine- but the second term is not! As a consequence, our Hilbert space is not appropriate to admit the action of Hamiltonian corresponding to the weak $|\Lambda\rangle$ decays.

Looking at the list of semi-leptonic decays of low-lying spin 1/2 baryons, one realizes that the same kind of pathology is present as soon as the state

$|abc\rangle$ is transformed into the state $|aac\rangle$, or inversely when the state $|aac\rangle$ is transformed into a combination of $|abc\rangle$ states, with a, b and c figuring out the quark flavours. But we see that this problem can easily be cured in the above considered $|\Lambda\rangle$ decay by choosing for $|\Lambda\rangle$ a wave function symmetric in u and s :

$$\Lambda = (u^\uparrow d^\downarrow s^\uparrow + s^\uparrow d^\downarrow u^\uparrow - u^\downarrow d^\uparrow s^\uparrow - s^\downarrow d^\uparrow u^\downarrow)/2$$

and in the general case by imposing the baryon wavefunctions to be completely symmetric in all the three quarks whether they were identical or not. The necessity of having a "complete" Hilbert space is what leads to the "generalized" Pauli principle.

Actually, it can be shown that the Hilbert space description in terms of completely symmetrized states in the three quarks is equivalent to the description in terms of states where the symmetry is imposed only for quarks of the same flavour. In other words, there is an isomorphism between the physics described in terms of two Hilbert spaces, the first one submitted to the usual Pauli principle and the second one submitted to what has been called the "generalized" Pauli principle. The mathematical proof is sketched in Appendix B.

Let us point out that, instead of the commonly adopted representation for the baryon wave-function (we call it "Generalized Pauli symmetric" in Appendix D), one might prefer the equivalent, simpler and more elegant expressions offered by the Fock space formalism, as developed in Appendix C and explicated in Appendix D for the octet of $S = 1/2$ low-lying baryons.

Finally, combining the completely antisymmetric colour part of the three quark wave functions with its completely symmetric complementary part provides, of course, with a wave function which is completely antisymmetric in the three constituent fermions, and that is namely what is prescribed by the generalized Pauli principle.

6 Conclusion

We have shown that the classification of baryons, and reasonable estimates of their masses, their magnetic moments, as well as of the corresponding

form factors relative to semi-leptonic decays can be obtained without any explicit reference to the flavour unitary group. All this can be included as well in fairly elementary quantum mechanics courses, as in elementary particle courses.

As widely explicated, the "minimal" or "partial" permutational symmetry imposed on wave-functions by the Pauli principle is well adapted as long as flavour changing forces do not operate. When they do, one is naturally led to the choice of wave functions satisfying complete symmetry under permutation of the three constituent quarks, that is, in other words, to the application of the generalized Pauli principle.

At this point it is interesting to make a connection between the above discussion in Section 5 and the classification of baryons in irreducible representations of the $SU(N)$ group, N being the number of different flavours. Actually, one knows that each irreducible representation \mathcal{R} of the $SU(N)$ group can be represented by a Young tableau with n boxes associated with an irreducible representation of the symmetric group \mathcal{S}_n also called permutation group \mathcal{P}_n of n elements. That means that any element of \mathcal{R} satisfies a special \mathcal{P}_n symmetry which is the same for all elements in \mathcal{R} . Taking as an example $SU(2)$, the representations above denoted $S=3/2$ and $S=1/2$ resulting from the tensorial product of three times the fundamental representation $S=1/2$ are respectively completely symmetric and "mixed symmetric" under \mathcal{P}_3 . Now, considering the group $SU(3)$, its irreducible eight dimensional 8 and ten dimensional 10 representations obtained by tensorial product of three fundamental representations 3 are respectively completely and mixed symmetric under \mathcal{P}_3 . Complete symmetry under flavour and spin will therefore be obtained by combining the spin- $SU(2)$ and flavour- $SU(3)$ representations as follows: $(S=1/2; \underline{8})$ and $(S=3/2; \underline{10})$. And we can here recognize the decomposition of the completely symmetric 56 representation of the so-called flavour-spin $SU(6)$ group commonly used to classify the S-wave baryons (in perfect accordance with results of Section 2). That also makes unitary groups so natural candidates for (broken) symmetry groups.

Acknowledgements :

P.S. is indebted to P. Aurenche and E. Guadagnini for numerous and valuable discussions. We both must specially thank R. Stora for his precious help in some algebraic aspects developed in this note, particularly in Appendices B and C, as well as for a critical reading of the manuscript.

Appendix A: Jacobi coordinates

The spatial wave functions ψ of our problem are the harmonic-oscillator eigenfunctions $\psi_{nLM}(\vec{\rho}, \vec{\lambda})$, where $\vec{\rho}$ and $\vec{\lambda}$ are the Jacobi coordinates which separate the Hamiltonian in Eq. (20) into two independent three-dimensional oscillators.

Hereafter, we give expressions for Jacobi coordinates that makes the separation of the center of mass motion in the case of unequal masses, and even when the interaction potential is different for each pair of particles.

So, starting from the Hamiltonian:

$$H_{\text{h.o.}} = \sum_i (m_i + p_i^2/2m_i) + \sum_{i<j} (k_{ij}/2) \cdot (\vec{r}_i - \vec{r}_j)^2 \quad (20)$$

with each coupling $k_{ij} = k_{ji}$ explicitly depending of the quark pair $q_i q_j$, one can rewrite the potential part as:

$$\sum_{i<j} (k_{ij}/2) \cdot (\vec{r}_i - \vec{r}_j)^2 = 1/2 \bar{Y} \mathcal{M} Y$$

with:

$$\bar{Y} = (\vec{y}_1, \vec{y}_2, \vec{y}_3)$$

and \vec{y}_i defined as: $\vec{y}_i = \sqrt{m_i} \vec{r}_i$, and the entries of the matrix \mathcal{M} satisfying:

$$\begin{aligned} \mathcal{M}_{ii} &= \frac{k_{ij} + k_{ik}}{m_i} \quad i \neq j \neq k \\ \mathcal{M}_{ij} &= \mathcal{M}_{ji} = -\frac{k_{ij}}{(m_i m_j)^{\frac{1}{2}}} \quad i \neq j \end{aligned}$$

The diagonalization of \mathcal{M} provides with one eigenvalue $\lambda_0 = 0$ associated to an eigenvector proportional to:

$$\vec{R}_{cm} = \frac{1}{M} (m_1 \vec{r}_1 + m_2 \vec{r}_2 + m_3 \vec{r}_3) \quad (21)$$

where: $M = m_1 + m_2 + m_3$, and two non zero eigenvalues λ_{\pm} respectively associated to the eigenvectors:

$$\vec{\rho} = h_+(m_1^{1/2}\vec{r}_1 + m_2^{1/2}f_+\vec{r}_2 + m_3^{1/2}g_+\vec{r}_3) \quad (22)$$

and:

$$\vec{\lambda} = h_-(m_1^{1/2}\vec{r}_1 + m_2^{1/2}f_-\vec{r}_2 + m_3^{1/2}g_-\vec{r}_3) \quad (23)$$

with:

$$h_{\pm} = \sqrt{\frac{\lambda_{\pm}}{(3K)^{\frac{1}{2}}(1 + f_{\pm}^2 + g_{\pm}^2)}} \\ K = k_{12}k_{13} + k_{21}k_{23} + k_{31}k_{32}$$

and:

$$f_{\pm} = \sqrt{\frac{m_2}{m_1}} \frac{[k_{23} + k_{31} - m_3\lambda_{\pm}][k_{31} + k_{12} - m_1\lambda_{\pm}] - k_{31}^2}{k_{23}k_{31} + k_{12}[k_{23} + k_{31} - m_3\lambda_{\pm}]} \\ g_{\pm} = \frac{\sqrt{m_3}}{k_{23}} \left[-\frac{k_{12}}{\sqrt{m_1}} + \frac{k_{23} + k_{12} - m_2\lambda_{\pm}}{\sqrt{m_2}} f_{\pm} \right]$$

Finally, the explicit expressions of λ_{\pm} read:

$$\lambda_{\pm} = 1/2 \left[\sum_{i; i \neq j \neq k \neq i} \frac{k_{ij} + k_{ik}}{m_i} \pm \sqrt{\left(\sum_{i; i \neq j \neq k \neq i} \frac{k_{ij} + k_{ik}}{m_i} \right)^2 - 4K \sum_{j < k} \frac{1}{m_j m_k}} \right]$$

In terms of \vec{R}_{cm} , $\vec{\rho}$, $\vec{\lambda}$, the total Hamiltonian acquires the following simple form in which the c.m. motion explicitly separates:

$$H_{h.o.} = \sum_i m_i + \frac{p_{\vec{R}_{cm}}^2}{2M} + \frac{p_{\vec{\rho}}^2}{2m_+} + \frac{p_{\vec{\lambda}}^2}{2m_-} + \frac{\sqrt{3K}}{2}(\vec{\rho}^2 + \vec{\lambda}^2) \quad (24)$$

where: $m_{\pm} = \frac{\sqrt{3K}}{\lambda_{\pm}}$

It might be useful to consider the special case with only one coupling constant, that is: $k_{ij} = k$ for $i, j = 1, 2, 3$. Then, K becomes $= 3k^2$ and, in the Hamiltonian, the potential part reduces to: $\frac{3}{2}k(\vec{\rho}^2 + \vec{\lambda}^2)$. In the kinetic part, one gets: $m_{\pm} = \frac{3k}{\lambda_{\pm}}$ and :

$$\lambda_{\pm} = k \left[\sum_i \frac{1}{m_i} \pm \sqrt{\sum_i \frac{1}{m_i^2} - \sum_{i < j} \frac{1}{m_i m_j}} \right]$$

while f_{\pm} and g_{\pm} simplify as :

$$f_{\pm} = \sqrt{\frac{m_2}{m_1}} \frac{4[1 - m_3(\lambda_{\pm}/2k)][1 - m_1(\lambda_{\pm}/2k)] - 1}{3 - 2m_3(\lambda_{\pm}/2k)}$$

$$g_{\pm} = \sqrt{m_3} \left[-\frac{1}{\sqrt{m_1}} + 2\frac{1 - m_2(\lambda_{\pm}/2k)}{\sqrt{m_2}} f_{\pm} \right]$$

Another relevant case is the one with two identical quarks, that is $m_1 = m_2 = m$ and $m_3 = m'$. Then it is reasonable to have : $k_{31} = k_{23} = k'$ and $k_{12} = k$, and the $\vec{\rho}$ and $\vec{\lambda}$ Jacobi vectors simply become:

$$\vec{\rho} = \frac{1}{\sqrt{2}} \left[\frac{2k + k'}{3k'} \right]^{\frac{1}{4}} (\vec{r}_1 - \vec{r}_2) \quad (25)$$

$$\vec{\lambda} = \frac{1}{\sqrt{2}} \left[\frac{k'}{3(2k + k')} \right]^{\frac{1}{4}} (\vec{r}_1 + \vec{r}_2 - 2\vec{r}_3) \quad (26)$$

Finally, we give the limit case with all three masses equal $m_1 = m_2 = m_3 = m$ and $k_{12} = k_{23} = k_{31} = k$:

$$\vec{\rho} = \frac{1}{\sqrt{2}} (\vec{r}_1 - \vec{r}_2) \quad (27)$$

$$\vec{\lambda} = \frac{1}{\sqrt{6}} (\vec{r}_1 + \vec{r}_2 - 2\vec{r}_3) \quad (28)$$

Appendix B: Generalized Pauli principle versus Pauli principle: a mathematical proof of their equivalence

This section is not directly dedicated to students. Of course, one can try and convince oneself on some examples that there is a one-to-one correspondence between the "mathematical" states and observables submitted to the Generalized Pauli principle (or GPP) and those simply satisfying the Pauli principle (or PP). The proof that we propose hereafter has two advantages. First, it is general, and so adaptable to situations other than that of regular baryons (for example multiquark states). Secondly, it introduces the Fock space formalism in a rather natural way. All that follows is taken from [3].

We start with some definitions :

Let k be a commutative ring, E a k -module, and N the set of integers.

Def.1: Symmetric tensor algebra:

Denoting $T^n(E)$ the set of elements: $z = x_1 \otimes x_2 \otimes \dots \otimes x_n$ with $x_1, \dots, x_n \in E$, the action of \mathcal{S}_n on $T^n(E)$ reads:

$$\sigma(x_1 \otimes x_2 \otimes \dots \otimes x_n) = x_{\sigma^{-1}(1)} \otimes x_{\sigma^{-1}(2)} \otimes \dots \otimes x_{\sigma^{-1}(n)} \quad (29)$$

The elements z such that:

$$\sigma \cdot z = z$$

are called symmetric tensors of order n . They form a sub- k -module of $T^n(E)$, denoted: $TS^n(E)$. One sets:

$$TS(E) = \bigoplus_{n=0}^{\infty} TS^n(E)$$

on which one can define a symmetric product (we do not give here the rule in order not to overload the text).

Def.2: Gamma algebra:

We call Gamma algebra of E and denote by $\Gamma(E)$ the associative, unifier, commutative, algebra defined by the set of generators $N \times E$ and relations:

$$(0, x) = 1 \quad (30)$$

$$(p, \lambda x) = \lambda^p (p, x) \quad (31)$$

$$(p, x + y) = \sum_{q=0}^p (p, x)(p - q, y) \quad (32)$$

Def.3: Exponential type sequence:

We call exponential type sequence of E elements a sequence $a = (a_p)_{p \in N}$ such that

$$a_0 = 1 \quad (33)$$

$$a_p a_q = \frac{(p+q)!}{p!q!} a_{p+q} \quad (34)$$

Now, we come to the results we need:

Prop.1: $\{(p, x)\}_{p \in N}$ with $x \in E$ is an exponential type sequence.

An interesting example of exponential type sequence is given by $f(x)$ with

$$f(x)_p = \frac{1}{p!} (x)^p \quad (35)$$

Prop.2: There exists one and only one isomorphism g between $\Gamma(E)$ and $TS(E)$:

$$\Gamma(E) \cong TS(E) \quad (36)$$

such that:

$$(g \circ \gamma_p)(x) = x \otimes x \otimes \dots \otimes x \quad (37)$$

with p x -factors, and where we have denoted by γ_p , with $p \in N$, the application from E into $\Gamma(E)$ product of the injection $x \rightarrow (p, x)$ and of the canonical homomorphism of the free commutative algebra $N \times E$ to $\Gamma(E)$.

Prop.3: Let E and F be two k modules. There exists one and only one isomorphism ϕ from $\Gamma(E \times F)$ into $\Gamma(E) \otimes_k \Gamma(F)$:

$$\Gamma(E \times F) \cong \Gamma(E) \otimes_k \Gamma(F) \quad (38)$$

such that:

$$(\phi \circ \gamma_p)(x, y) = \sum_{q=0}^p \gamma_q(x) \otimes \gamma_{p-q}(y) \quad (39)$$

The above properties can be reformulated by considering the case of polynomials. Indeed, one directly remarks that $\Gamma(E)$ can be seen as the set of polynomials on the dual E^* of E , that is also:

$$\Gamma(E) \cong \text{Polyn}(E^*) \quad (40)$$

Now, considering E finite dimensional, that is made of elements $\vec{x} = \sum_i x^i e_i \in E$ where $(e_i)_{i \in I}$ is a basis of E , with $x_i \in k$ for each $i \in I$, and I being finite,

then Prop.2 insures that:

$$\Gamma(E) \cong \text{Polyn}(E^*) \cong (\text{Polyn}(E))^* \cong TS(E) \quad (41)$$

and from Prop.3 we get:

$$\text{Polyn}(\vec{x}, \vec{y}) = \text{Polyn}(\vec{x}) \otimes \text{Polyn}(\vec{y}) \quad (42)$$

with $\vec{x} \in E$ and $\vec{y} \in F$, keeping in mind the relations (31),(34) and (38).

Coming back to our flavoured quarks, let us associate to each flavour $f = u, d, \dots$ the $S = 1/2$ spin representation space E_f generated by f^\uparrow, f^\downarrow . One can write:

$$\Gamma(E_f) \cong \text{Polyn}(f^\uparrow, f^\downarrow) \quad (43)$$

and also :

$$\Gamma(E_f \times E_{f'}) \cong \Gamma(E_f) \otimes_R \Gamma(E_{f'}) \quad (44)$$

our states being defined on the ring of real numbers R .

Appendix C: Fock space wave-functions

Actually, the above developed framework is nothing else than Fock space. Let us associate to each just above defined couple f^\uparrow, f^\downarrow in E_f the - commuting or "bosonic" - creator operators $a_{f^\uparrow}^\dagger, a_{f^\downarrow}^\dagger$, acting on the $|0\rangle$ vacuum. Again, we can write:

$$\Gamma(E_f) \cong \text{Polyn}(a_{f^\uparrow}^\dagger, a_{f^\downarrow}^\dagger) \quad (45)$$

and also:

$$\text{Polyn}(a_{f^\uparrow}^\dagger, a_{f^\downarrow}^\dagger) \otimes \text{Polyn}(a_{f'^\uparrow}^\dagger, a_{f'^\downarrow}^\dagger) \cong \text{Polyn}(a_{f^\uparrow}^\dagger, a_{f^\downarrow}^\dagger, a_{f'^\uparrow}^\dagger, a_{f'^\downarrow}^\dagger) \quad (46)$$

Due to the invariance under permutations of a monomial in a_f^\dagger operators, one can get, for baryon states, expressions as simple as the ones used in Section (4) and denoted in Appendix D below "simply Pauli symmetric" - so much simpler than those called "generalized Pauli symmetric" in the same appendix - but unambiguously defined, as these last ones, for computations involving flavour changes. More precisely, it is rather immediate to

translate a wave-function written in our "Pauli symmetric basis" into the corresponding one in the Fock basis. One has just to take care of the respective normalization of states: as an example, the spin $S = 1, S_z = 0$ state $(u^\uparrow u^\downarrow + u^\downarrow u^\uparrow)/\sqrt{2}$ will correspond to $a_{u^\uparrow}^\dagger a_{u^\downarrow}^\dagger$, while $u^\uparrow u^\uparrow$ will be associated to $a_{u^\uparrow}^\dagger a_{u^\uparrow}^\dagger/\sqrt{2}$, using the usual normalized power product $(a_f^\dagger)^n/\sqrt{n!}$. An explicit comparison of the wave-functions for $S = 1/2$ low-lying baryons in the different bases can be performed with Appendix D.

Finally, associating, as usual, to each creator a_f^\dagger its annihilator counter part a_f such that $a_f|0\rangle = 0$ and $(a_f^\dagger)^\dagger = a_f$, and satisfying the well-known Heisenberg algebra commutation relations:

$$[a_i, a_j^\dagger] = \delta_{ij} \quad (47)$$

$$[a_i, a_j] = [a_i^\dagger, a_j^\dagger] = 0 \quad (48)$$

with i, j representing any f^\uparrow or f^\downarrow , simple expressions can be given for the self-adjoint parts of the Hamiltonian representing the flavour changing f into f' :

$$H_{f \rightarrow f'}^\pm = (a_{f'^\uparrow}^\dagger a_{f^\uparrow} \pm a_{f'^\downarrow}^\dagger a_{f^\downarrow} + h.c.) \quad (49)$$

involved in the computations of the vector f_1 and axial g_1 form factors of baryon semi-leptonic decays.

It might look surprising to the reader that we use a bosonic Fock space although we are dealing with fermionic quarks. We should not forget that antisymmetry is carried by the colour part. In other words, let us say that in the fermionic space of coloured quarks, there exists a subspace constituted by three quark states which are colour singlets - our baryons - this subspace being isomorphic, as a Hilbert space, to the three particles subspace of a bosonic Fock space.

Appendix D: Wave-functions of spin 1/2 S-wave baryons

Different expressions for the wave-functions of baryons have been used following the kind of computations we had to consider. As long as classification

and flavour non changing quantities are involved, wave-functions satisfying the symmetry imposed by the Pauli principle are perfectly adapted: we will call them "partial symmetric" or "simple Pauli symmetric" wave-functions. As soon as flavour changing forces are involved, complete permutational invariance of the three quarks is required : we will denoted them "generalized Pauli symmetric" wave-functions. In this last case, the use of the Fock space formalism, as seen in Appendix C, looks rather appealing: their natural denomination will be "Fock" wave-functions. Let us once more mention that the colour part, and so the antisymmetry nature of the baryon, has been "factored out", that is why one is concerned below only with permutational symmetry aspects, and in particular with bosonic Fock operators.

Simple Pauli symmetric wave-functions

$$N = [2d^\uparrow d^\uparrow u^\downarrow - (d^\uparrow d^\downarrow + d^\downarrow d^\uparrow)u^\uparrow]/\sqrt{6}$$

$$P = [2u^\uparrow u^\uparrow d^\downarrow - (u^\uparrow u^\downarrow + u^\downarrow u^\uparrow)d^\uparrow]/\sqrt{6}$$

$$\Lambda = [(u^\uparrow d^\downarrow - u^\downarrow d^\uparrow)s^\uparrow]/\sqrt{2}$$

$$\Sigma^O = [2u^\uparrow d^\uparrow s^\downarrow - (u^\uparrow d^\downarrow + u^\downarrow d^\uparrow)s^\uparrow]/\sqrt{6}$$

$$\Sigma^+ = [2u^\uparrow d^\uparrow s^\downarrow - (u^\uparrow u^\downarrow + u^\downarrow u^\uparrow)s^\uparrow]/\sqrt{6}$$

$$\Sigma^+ = [2d^\uparrow d^\uparrow s^\downarrow - (d^\uparrow d^\downarrow + d^\downarrow d^\uparrow)s^\uparrow]/\sqrt{6}$$

$$\Xi^- = [2s^\uparrow s^\uparrow d^\downarrow - (s^\uparrow s^\downarrow + s^\downarrow s^\uparrow)d^\uparrow]/\sqrt{6}$$

$$\Xi^0 = [2s^\uparrow s^\uparrow u^\downarrow - (s^\uparrow s^\downarrow + s^\downarrow s^\uparrow)u^\uparrow]/\sqrt{6}$$

Generalized Pauli symmetric wave-functions

$$N = [2d^\uparrow d^\uparrow u^\downarrow - (d^\uparrow d^\downarrow + d^\downarrow d^\uparrow)u^\uparrow + \text{permutations}]/\sqrt{18}$$

$$P = [2u^\uparrow u^\uparrow d^\downarrow - (u^\uparrow u^\downarrow + u^\downarrow u^\uparrow)d^\uparrow + \text{permutations}]/\sqrt{18}$$

$$\Lambda = [(u^\uparrow d^\downarrow - u^\downarrow d^\uparrow)s^\uparrow + \text{permutations}]/\sqrt{12}$$

$$\Sigma^O = [2u^\uparrow d^\uparrow s^\downarrow - (u^\uparrow d^\downarrow + u^\downarrow d^\uparrow)s^\uparrow + \text{permutations}]/6$$

$$\Sigma^+ = [2u^\uparrow d^\uparrow s^\downarrow - (u^\uparrow u^\downarrow + u^\downarrow u^\uparrow)s^\uparrow + \text{permutations}]/\sqrt{18}$$

$$\Sigma^- = [2d^\uparrow d^\uparrow s^\downarrow - (d^\uparrow d^\downarrow + d^\downarrow d^\uparrow)s^\uparrow + \text{permutations}]/\sqrt{18}$$

$$\Xi^- = [2s^\uparrow s^\uparrow d^\downarrow - (s^\uparrow s^\downarrow + s^\downarrow s^\uparrow)d^\uparrow + \text{permutations}]/\sqrt{18}$$

$$\Xi^0 = [2s^\uparrow s^\uparrow u^\downarrow - (s^\uparrow s^\downarrow + s^\downarrow s^\uparrow)u^\uparrow + \text{permutations}]/\sqrt{18}$$

Fock wave-functions

$$N = \frac{1}{\sqrt{3}} [a_{d^\uparrow}^\dagger a_{d^\uparrow}^\dagger a_{u^\downarrow}^\dagger - a_{d^\uparrow}^\dagger a_{d^\downarrow}^\dagger a_{u^\uparrow}^\dagger] |0\rangle$$

$$P = \frac{1}{\sqrt{3}} [a_{u^\uparrow}^\dagger a_{u^\uparrow}^\dagger a_{d^\downarrow}^\dagger - a_{u^\uparrow}^\dagger a_{u^\downarrow}^\dagger a_{d^\uparrow}^\dagger] |0\rangle$$

$$\Lambda = \frac{1}{\sqrt{2}} [a_{u^\uparrow}^\dagger a_{d^\downarrow}^\dagger a_{s^\uparrow}^\dagger - a_{u^\downarrow}^\dagger a_{d^\uparrow}^\dagger a_{s^\uparrow}^\dagger] |0\rangle$$

$$\Sigma^0 = \frac{1}{\sqrt{6}} [2a_{u^\uparrow}^\dagger a_{d^\uparrow}^\dagger a_{s^\downarrow}^\dagger - a_{u^\uparrow}^\dagger a_{d^\downarrow}^\dagger a_{s^\uparrow}^\dagger - a_{u^\downarrow}^\dagger a_{d^\uparrow}^\dagger a_{s^\uparrow}^\dagger] |0\rangle$$

$$\Sigma^+ = \frac{1}{\sqrt{3}} [a_{u^\uparrow}^\dagger a_{u^\uparrow}^\dagger a_{s^\downarrow}^\dagger - a_{u^\uparrow}^\dagger a_{u^\downarrow}^\dagger a_{s^\uparrow}^\dagger] |0\rangle$$

$$\Sigma^- = \frac{1}{\sqrt{3}} [a_{d^\uparrow}^\dagger a_{d^\uparrow}^\dagger a_{s^\downarrow}^\dagger - a_{d^\uparrow}^\dagger a_{d^\downarrow}^\dagger a_{s^\uparrow}^\dagger] |0\rangle$$

$$\Xi^- = \frac{1}{\sqrt{3}} [a_{s^\uparrow}^\dagger a_{s^\uparrow}^\dagger a_{d^\downarrow}^\dagger - a_{s^\uparrow}^\dagger a_{s^\downarrow}^\dagger a_{d^\uparrow}^\dagger] |0\rangle$$

$$\Xi^0 = \frac{1}{\sqrt{3}} [a_{s^\uparrow}^\dagger a_{s^\uparrow}^\dagger a_{u^\downarrow}^\dagger - a_{s^\uparrow}^\dagger a_{s^\downarrow}^\dagger a_{u^\uparrow}^\dagger] |0\rangle$$

References

- [1] B.R.Martin and G.Shaw : *Particle Physics*, John Wiley and Sons, Chichester(1997)
- [2] J. Franklin, Phys. Rev. **172** (1968) 1807.
- [3] N. Bourbaki, *Eléments de Mathématique, Algèbre*, Chap. 4-7, Masson Ed.(1981), p.87-89.